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DECAYS OF HIGGS SCALARS INTO VECTOR MESONS AND PHOTONS ☆

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Decays of a neutral Higgs scalar (H) into a neutral vector meson (V) and a photon (γ) are discussed. The width for $H \to V + \gamma$ is related to the leptonic width $V \to l^+l^-$. The rate for $H \to V + V$ is also estimated using a phenomenological quark interaction with vector mesons supplemented by form factors.

Higgs scalars are an important ingredient in the construction of unified gauge theories of weak and electromagnetic interactions. Recent results from a wide variety of experiments support the simple SU(2) $\times$ U(1) gauge theory constructed by Weinberg and Salam [1]. In principle one can develop unified models which are phenomenological in nature and which are also consistent with experiments. Higgs scalars have no place in such alternatives to gauge models [2]. Discovery of Higgs scalars with prescribed couplings would therefore constitute a most important verification of spontaneously broken gauge theories. The mass of the Higgs scalar, being a free parameter in the theory [1], may be less than the mass of the W or Z boson and therefore the Higgs meson may well be accessible, while the W or Z bosons are not, to existing machines or those that are to come in operation soon. As a result, considerable effort is being directed towards understanding the phenomenology of these particles [3–7].

Detection of Higgs scalars is likely to be a formidable experimental problem. Their coupling to fermions (f of mass m_f) $g_f = G_F^{1/2}\, 2^{1/4} m_f$ suggests that their decays to τ (or heavier) leptons would be important and perhaps relatively clean ‡1. However, in practice the τ leptons from Higgs decay may not be appreciable compared to the "background" production of $\tau^+\tau^-$. For example, in pp collision where the Higgs cross section is appreciable [6] the rate for $pp \to \tau^+\tau^- X$ through virtual photons [8] is many times more than the τ's from $pp \to H + X$ followed by $H \to \tau^+\tau^-$. Missing-mass techniques [4,5,7] may have the best chance of revealing an H. Its detailed properties may have to be studied for confirmation that it is indeed a Higgs particle. Two-body decays and/or decay modes in which all the products are detectable could be particularly useful in this regard. One is thus led to consider neutral Higgs decays into $\gamma + \gamma$, $\gamma + V$, $V + V$, where γ is the photon and V a neutral vector meson. Unfortunately $H \to \gamma + \gamma$ turns out to be small because of a subtle cancellation arising from the contributions of fermions and weak boson loops [3]. Such a cancellation can be avoided if one replaces one or both photons with strongly interacting vector mesons. We are thus motivated to investigate $H \to V + \gamma$ and $H \to V + V$.

These decays proceed through the quark loop as shown in fig. 1. We will concern ourselves with the case where the vector meson (V) is a pure orthoquarkonium bound state of the particular quark flavor (of charge $|e_q|$) in the loop. For $m_H \sim m_V$ one can calculate the width $\Gamma_{HV\gamma}$ for the reaction $H \to V + \gamma$ in terms of $|\psi(0)|^2$, where $\psi(0)$ is the radial wave function of the V at the origin ‡2:

$$\Gamma^{b}_{HV\gamma} = 48 g_f^2 \alpha e_q^2 |\psi(0)|^2 (1-r)/m_H^2(1+r) \tag{1}$$

$$= 3 g_f^2 r(1-r)\Gamma_{V\ell\ell}/\alpha\pi(1+r), \tag{2}$$

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‡1 For definiteness we will concern ourselves with the simplest spontaneously broken gauge theory of Weinberg–Salam [1] which involves a single neutral physical Higgs boson.

‡2 Superscripts b and g in eqs. (1) and (3) are used to distinguish the bound state calculation from the perturbative one.

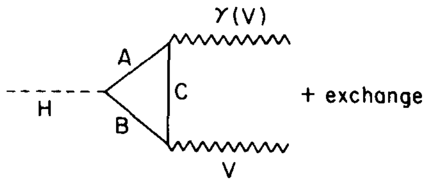

Fig. 1. Diagrams for $H \to V + \gamma(V)$. H is Higgs, A, B, C are internal lines of a particular quark flavor that V is composed of.

where $r = m_V^2/m_H^2$ and $\Gamma_{V\ell\ell}$ is the leptonic width $V \to ee$ or $V \to \mu\mu$ [9,10]. This calculation for the width of $H \to V + \gamma$ in terms of $\psi(0)$ corresponds to taking the quark lines B and C on-shell. As m_H grows larger, the contribution from the real part of the loop becomes increasingly important and thus $\Gamma^g_{HV_\gamma}$ would be expected to be a lower bound to the actual width for $m_H \gg m_V$.

For calculating Γ_{HV_γ} for $m_H \gg m_V$ we imagine that the vector meson has an effective interaction with quarks of the form $-ig_V \bar{\psi}_q \gamma_\mu \psi_q V^\mu$. Setting $2m_q = m_V$ we are led to:

$$\Gamma^g_{HV_\gamma} = 9 e_q^2 \alpha \alpha_V g_f^2 m_H r(1-r)|I_\gamma(r)|^2/128\pi^3 , \qquad (3)$$

where $\alpha_V = g_V^2/4\pi$ and $I_\gamma(r)$ is the loop integral:

$$I_\gamma(r) = \int_0^1 dx \int_0^x dy [1 - 4y(1-x)] \times [\tfrac{1}{4} r(2x-1)^2/1 - r - y(1-x) - i\epsilon]^{-1} . \qquad (4)$$

Several remarks are in order. First we will content ourselves with six quark flavors with the sixth t quark of charge $e_q = +2/3$ [‡3]. Secondly, we will, following Sakurai's observation [12], assume that $\Gamma_{V\ell\ell}/e_q^2 \approx 1.3$ keV is a constant for all quark flavors. This constancy with the assumed quark–vector meson interaction implies (from dimensional considerations) that $\alpha_V \propto 1/m_V$. The constant of proportionality is fixed by normalizing the width (3) with the bound state result (2) in a region where the latter is expected to hold best, that is for $m_H \sim m_V$. For definiteness, we did the normalization at $m_H = \sqrt{2}\, m_\psi$ for $H \to \psi + \gamma$. In this way we find:

$$\alpha_V \approx 3.9 \text{ GeV}/m_V . \qquad (5)$$

[‡3] We assume that $\Upsilon(9.46)$ is composed of b quarks of charge $-1/3$. See ref. [11].

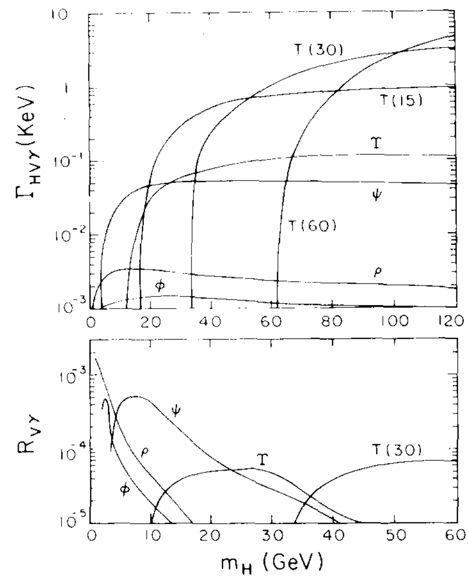

Fig. 2. Decay width Γ_{HV_γ} for $H \to V + \gamma$ and R_{V_γ} as defined by eq. (7) are shown for various vector mesons as a function of m_H.

Numerical results for $H \to V + \gamma$ are presented in fig. 2. The figure shows the partial width as a function of m_H for ρ, φ, ψ, Υ and for a hypothetical heavier vector meson T made of t quark–antiquark with masses m_T = 15, 30, or 60 GeV. The curves rise sharply above threshold and attain a constant value for $m_H/m_V \gtrsim 4$.

From the figure we notice that for $m_H^2/m_V^2 \gg 1$ a rough estimate for the width Γ_{HV_γ} can be obtained by using the following simple relation [‡4]

$$\Gamma_{HV_\gamma}(m_H^2/m_V^2 \gg 1) \approx 10^{-2} e_q^2 (m_V/\text{GeV})^2 \text{ keV} . \qquad (6)$$

Of greater experimental interest is the branching ratio $\Gamma(H \to V + \gamma)/\Gamma(H \to \text{all})$, whose denominator is a complicated function of m_H. To compute it reliably one needs to know all the thresholds for a given m_H. So we have computed what we believe is the next best thing:

$$R_{VA} = \Gamma(H \to V + A)/\sum_f \Gamma(H \to f\bar{f}) , \qquad (7)$$

where $A = \gamma$ for $H \to V + \gamma$ and Σ stands for the sum

[‡4] This formula holds for the range of masses shown in the figure but is inapplicable if m_H/m_V is astronomically large.

over the leptons e, μ, τ and quark flavors u, d, s, c, b and t and

$$\Gamma(\mathrm{H} \to \mathrm{f}\bar{\mathrm{f}}) = C_{\mathrm{f}} g_{\mathrm{f}}^2 m_{\mathrm{H}} [1 - 4m_{\mathrm{f}}^2/m_{\mathrm{H}}^2]^{3/2}/8\pi\,, \tag{8}$$

where $C_{\mathrm{f}} = 3$ for quarks and $C_{\mathrm{f}} = 1$ for leptons. R_{VA} is thus the upper bound on the branching ratio for $\mathrm{H} \to \mathrm{V} + \mathrm{A}$ [‡5]. Fig. 2 shows $R_{\mathrm{V}\gamma}$ for various vector mesons as a function of m_{H}. Notice that for $m_{\mathrm{H}} \lesssim 4$ GeV, $\mathrm{H} \to \rho + \gamma$ is dominant, for 4 GeV $\lesssim m_{\mathrm{H}} \lesssim 25$ GeV ψ dominates, for $m_{\mathrm{H}} \gtrsim 25$ GeV decays into Υ are slightly more than into ψ and with $m_{\mathrm{T}} = 30$ GeV decays $\mathrm{H} \to \mathrm{T} + \gamma$ are dominant for $m_{\mathrm{H}} \gtrsim 40$ GeV.

For $m_{\mathrm{H}} \sim m_{\mathrm{V}}$ the bound state result (2) for $\mathrm{H} \to \mathrm{V} + \gamma$, in terms of the leptonic width of the vector mesons is quite reliable. For $m_{\mathrm{H}} \gg m_{\mathrm{V}}$ the partial width $\Gamma_{\mathrm{HV}\gamma}$ depends somewhat (but not too appreciably) on the value of $m_{\mathrm{H}}/m_{\mathrm{V}}$ that one chooses to normalize α_{V} at. However, the ratio of decay rates for $\mathrm{H} \to \mathrm{V} + \gamma$ into any two vector mesons is independent of the normalization for α_{V} and is, therefore, reliable. The decays $\mathrm{H} \to \mathrm{V} + \mathrm{V}$, on the other hand, *cannot* be calculated accurately. These modes may be quite important and even rough estimates for them (which can be obtained rather easily from our previous calculation) could be of experimental use. For this purpose we have adopted a purely phenomenological approach and estimated the partial width Γ_{HVV} by supplementing the effective quark–vector meson interaction with form factors $F_n(q_{\mathrm{mn}}^2)$ leading to

$$\Gamma_{\mathrm{HVV}} = 9\alpha_{\mathrm{V}}^2 g_{\mathrm{f}}^2 m_{\mathrm{H}} (1 - 4r)^{1/2} r [1 + 2r^2/(1 - 2r)^2] \times |I_{\mathrm{V}}(r)|^2 F_n(q_{\mathrm{mn}}^2)/256\pi^3\,, \tag{9}$$

where

$$I_{\mathrm{V}}(r) = \int_0^1 \mathrm{d}x \int_0^x \mathrm{d}y\, [1 - 4y(1-x)] \times \{[r/4(1-2r)]\,[(2x-1)^2 - 4y(1-y)] - y(1-x) - \mathrm{i}\epsilon\}^{-1}\,, \tag{10}$$

$$F_n(q_{\mathrm{mn}}^2) = [1 + (m_{\mathrm{q}}^2 - q_{\mathrm{mn}}^2)/m_{\mathrm{q}}^2]^{-n}\,. \tag{11}$$

Here q_{mn}^2 is the minimum of the squared momentum of quark line C (fig. 1) when lines A and B are on their mass shell and is given by

$$q_{\mathrm{mn}}^2 = 1.25 m_{\mathrm{V}}^2 - 0.5 m_{\mathrm{H}}^2 [1 - (1-r)^{1/2}(1-4r)^{1/2}]\,. \tag{12}$$

It may be worth pointing out that q_{mn}^2 is timelike for $\mathrm{H} \to \mathrm{V} + \gamma$ and space-like for $\mathrm{H} \to \mathrm{V} + \mathrm{V}$. Now as line C goes away from mass shell F_n (with $n > 1$) would tend to bring about the suppression of the decay mode. For $m_{\mathrm{H}} \sim m_{\mathrm{V}}$, q_{mn}^2 is $\sim -3m_{\mathrm{q}}^2$ and for $m_{\mathrm{H}} \gg m_{\mathrm{V}}$, $q_{\mathrm{mn}}^2 \sim -9m_{\mathrm{q}}^4/m_{\mathrm{H}}^2$, so that near threshold $F_n \to 5^{-n}$ and asymptotically $F_n \to 2^{-n}$. For $n = 2$ the resulting suppression ranges from 1/25 to 1/4. It may be worth mentioning that since $\alpha_{\mathrm{V}} \approx 3.9$ GeV$/m_{\mathrm{V}}$, for $m_{\mathrm{V}} \gtrsim 4$ GeV, $\alpha_{\mathrm{V}} \lesssim 1$, consequently for large values of m_{V} a perturbative calculation for $\mathrm{H} \to \mathrm{V} + \mathrm{V}$ may turn out to be reasonable anyway. Fig. 3 shows the numerical results for the partial width (Γ_{HVV}) and R_{VV}, defined by eq. (7) with A = V for ρ, ψ, Υ, and T with a mass of 30 GeV. For $m_{\mathrm{H}} < 8$ GeV decays into ρ, ω, φ are dominant, for $10 \lesssim m_{\mathrm{H}} \lesssim 28$ GeV decays into a pair

[‡5] For numerical computations we have assumed quark mass $\approx \frac{1}{2} \times$ the mass of the corresponding vector meson. In practice, the effective quark mass (m_{q}) for the decay $\mathrm{H} \to \mathrm{q}\bar{\mathrm{q}}$ may be somewhat more. This could appreciably change the value of $\Sigma\Gamma(\mathrm{H} \to \mathrm{f}\bar{\mathrm{f}})$ and consequently the ratio R_{VA} in the range $m_{\mathrm{V}} < m_{\mathrm{H}} < 2m_{\mathrm{q}}$.

[‡6] The decay $\mathrm{H} \to \mathrm{V} + \mathrm{V}$ depends on the quark mass and not on its charge. ϕ and ω are, therefore, omitted to avoid overcrowding fig. 3 as their curves are not appreciably different from that for ρ.

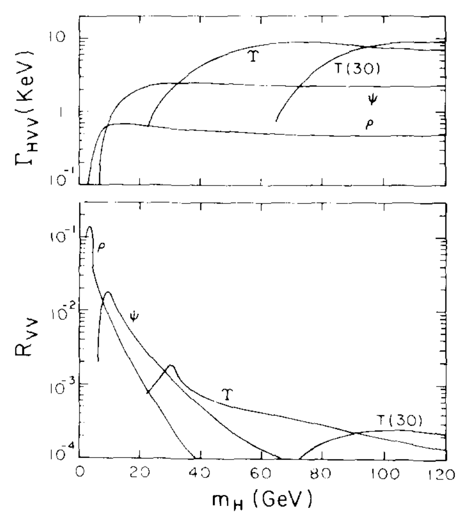

Fig. 3. Decay width Γ_{HVV} for $\mathrm{H} \to \mathrm{V} + \mathrm{V}$ and R_{VV} defined by eq. (7) are shown for various vector mesons as a function of M_{H} [‡6].

of ψ dominate and for $m_H \gtrsim 30$ GeV, $H \to \Upsilon + \Upsilon$ is the dominant mode.

Let us finally discuss, in brief, how these decay modes and this calculation may be put to use. First consider the V + V mode. Here one possible final signature for a Higgs could be $\ell^+\ell^- X$ via (ℓ = e or u):

$$\begin{array}{l} H \to V + V, \\ \quad | \quad \llcorner\!\!\to \ell^+\ell^- \\ \quad \llcorner\!\!\to X(\text{no } \ell) \end{array} \tag{13}$$

where the X contains no ℓ. The final signal $\ell^+\ell^- X$ would be additionally suppressed by the leptonic branching ratio of V. From fig. 3 we see that the rates for $H \to V + V$ (at least for $m_H \lesssim 30$ GeV and perhaps even for larger m_H) may be large enough to allow its experimental observation through reaction (13).

Next consider $H \to V + \gamma$. In principle, one would like to see the V through its leptonic decay mode and then search for a peak in $\ell^+\ell^-\gamma$ mass distribution where $m_{\ell^+\ell^-} = m_V$. The numbers in fig. 2 are already small enough that, by the time one folds in the leptonic branching ratio of the V, searching for a Higgs peak in $\ell^+\ell^-\gamma$ may become almost impossible. Somewhat more realistic, but still difficult, may be the search for γ rays resulting from a two-body decay $H \to \gamma + V$ followed by $V \to X$. That is, one has to select γ + x events such that $m_x = m_V$.

The most likely role that these decay modes may play lies in confirming or refuting that a given candidate discovered through, say, a missing-mass technique is a Higgs scalar or not. The point is that any spin-zero boson can in principle decay to a $V + \gamma$ or $V + V$. Since the theoretical uncertainty in $H \to V + \gamma$, especially in the ratio of the rates of the decays into two different vector mesons, is minimal, if the measured rates of a given candidate for Higgs are (say) too large compared to theoretical expectations, then it simply cannot be a Higgs scalar.

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